

# Quasifree kaon-photoproduction from nuclei in a relativistic approach

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## Abstract

We compute the recoil polarization of the lambda-hyperon and the photon asymmetry for the quasifree photoproduction of kaons in a relativistic impulse-approximation approach. Our motivation for studying polarization observables is threefold. First, polarization observables are more effective discriminators of subtle dynamics than the unpolarized cross section. Second, earlier nonrelativistic calculations suggest an almost complete insensitivity of polarization observables to distortions effects. Finally, this insensitivity entails an enormous simplification in the theoretical treatment. Indeed, by introducing the notion of a “bound-nucleon propagator” we exploit Feynman’s trace techniques to develop closed-form, analytic expressions for all photoproduction observables. Moreover, our results indicate that polarization observables are also insensitive to relativistic effects and to the nuclear target. Yet, they are sensitive to the model parameters, making them ideal tools for the study of modifications to the elementary amplitude — such as in the production, propagation, and decay of nucleon resonances — in the nuclear medium.

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## I. INTRODUCTION

Impelled by recent experimental advances, there is an increasing interest in the study of strangeness-production reactions from nuclei. These reactions form our gate to the relatively unexplored territory of hypernuclear physics. Moreover, these reactions constitute the basis for studying novel physical phenomena, such as the existence of a kaon condensate in the interior of neutron stars [1]. Indeed, the possible formation of the condensate could be examined indirectly by one of the approved experiments [2] at the Thomas Jefferson National Accelerator Facility (TJNAF). This experimental approach is reminiscent of the program carried out at the Los Alamos Meson Physics Facility (LAMPF) where pion-like modes were studied extensively through the quasifree ( $\vec{p}, \vec{n}$ ) reaction [3,4]. These measurements placed strong constraints on the (pion-like) spin-longitudinal response and showed conclusively that the long-sought pion-condensed state does not exist.

The work presented here is a small initial step towards a more ambitious program that concentrates on relativistic studies of strangeness in nuclei. Our aim in this paper is the study of the photoproduction of kaons from nuclei in the quasifree regime. This investigation helps us in two fronts. First, it sheds light on the elementary process,  $\gamma p \rightarrow K^+ \Lambda$ , by providing a different physical setting (away from the on-shell point) for studying the elementary amplitude. Second, it will enable us, in a future study, to explore modifications to the kaon propagator in the nuclear medium and to search for those observables most sensitive to the formation of the condensate. To achieve these goals we focus on the study of polarization observables. Polarization observables have been instrumental in the understanding of elusive details about subatomic interactions, as they become much more effective discriminators of subtle physical effects than the traditional unpolarized cross section. Moreover, quasifree polarization observables might be one of the cleanest tools for probing the nuclear dynamics. For example, the reactive content of the process is simple, as it is dominated by the quasifree production and knockout of a  $\Lambda$ -hyperon. Further, free polarization observables provide a baseline, against which possible medium effects may be inferred. Deviations of polarization observables from their free values are likely to arise from a modification of the interaction inside the nuclear medium or from a change in the response of the target. Indeed, relativistic models of nuclear structure predict medium modifications to the free observables stemming from an enhanced lower component of the Dirac spinors in the nuclear medium [5]. Finally, nonrelativistic calculations of the photoproduction of pseudoscalar mesons suggest that, while distortion effects provide an overall reduction of the cross section, they do so without substantially affecting the shape of the distribution [6–8]. Indeed, these nonrelativistic calculations show that two important polarization observables — the recoil polarization of the ejected baryon and the photon asymmetry — are insensitive to distortion effects. Moreover, they seem to be also independent of the mass of the target nucleus. Thus, quasifree polarization observables might represent a fundamental property of nuclear matter.

The insensitivity of polarization observables to distortion effects is clearly of enormous significance, as one can unravel distortion effects from those effects arising from relativity or from the large-momentum components in the wavefunction of the bound nucleon. Indeed, relativistic plane-wave impulse approximation (RPWIA) calculations have been successful in identifying physics not present at the nonrelativistic level. For example, relativistic effects

have been shown to contaminate any attempt to infer color transparency from a measurement of the asymmetry in the  $(e, e'p)$  reaction [9]. Further, the well-known factorization limit of nonrelativistic plane-wave calculations has been shown to break down due to the presence of negative-energy components in the bound-nucleon wavefunction [10]. Finally, neglecting distortions affords the computation of all polarization observables in closed form [9] by using the full power of Feynman's trace techniques.

We have organized our paper as follows. In Sec. II we discuss in detail our relativistic plane-wave formalism, placing special emphasis on the “bound-nucleon propagator” and on the use of Feynman's trace techniques to evaluate all observables. Our results are presented in Sec. III, where polarization observables are computed with two alternative parameterizations of the elementary amplitude and then compared to results obtained from an on-shell nucleon. We offer our conclusions and perspectives for future work in Sec. IV.

## II. FORMALISM

The kinematics for the quasifree production of a  $\Lambda$ -hyperon through the photoproduction reaction  $A(\gamma, K^+\Lambda)B$  is constrained by two conditions. First, there is an overall energy-momentum conservation:

$$k + p_A = k' + p' + p_B . \quad (1)$$

Note that  $k$  is the four-momentum of the incident photon, while  $k'$  and  $p'$  are the momenta of the produced kaon and lambda-hyperon, respectively. Finally,  $p_A(p_B)$  represents the momentum of the target(residual) nucleus. Moreover, since we are studying the photoproduction process within the framework of the impulse approximation (see Fig. 1) there is a second kinematical constraint arising from energy-momentum conservation at the  $\gamma N \rightarrow K^+\Lambda$  vertex:

$$k + p = k' + p' , \quad (2)$$

where  $p$  is the four-momentum of the bound nucleon, whose space part is known as the missing momentum:

$$\mathbf{p}_m \equiv \mathbf{p}' - \mathbf{q} ; \quad (\mathbf{q} \equiv \mathbf{k} - \mathbf{k}') . \quad (3)$$

Thus, as in most semi-inclusive processes — such as in the  $(e, e'p)$  reaction — the quasifree production process becomes sensitive to the nucleon momentum distribution. The differential cross section for the quasifree process is derived to be:

$$\left( \frac{d^5\sigma(s', \varepsilon)}{d\mathbf{k}' d\Omega_{\mathbf{k}'} d\Omega_{\mathbf{p}'}} \right)_{\text{lab}} = \frac{2\pi}{2E_\gamma} \frac{|\mathbf{k}'|^2}{(2\pi)^3 2E_{\mathbf{k}'}} \frac{M_N |\mathbf{p}'|}{(2\pi)^3} |\mathcal{M}|^2 , \quad (4)$$

where  $s'$  is the spin of the emitted  $\Lambda$ ,  $\varepsilon$  is the polarization of the incident photon,  $M_N$  is the nucleon mass, and  $\mathcal{M}$  is the transition matrix element given by:

$$|\mathcal{M}|^2 = \sum_m \left| \bar{U}(\mathbf{p}', s') T(s, t) U_{\alpha, m}(\mathbf{p}) \right|^2 . \quad (5)$$

Here  $\mathcal{U}(\mathbf{p}', s')$  is the free Dirac spinor for the emitted  $\Lambda$ -hyperon and  $\mathcal{U}_{\alpha,m}(\mathbf{p})$  is the Fourier transform of the relativistic spinor for the bound nucleon ( $\alpha$  denotes the collection of all quantum numbers necessary to specify the single-particle orbital). Note that we assume the impulse approximation valid and employ the on-shell photoproduction operator  $T(s, t)$ .

The unpolarized differential expression can be obtained by summing over the two possible components of the spin of the  $\Lambda$  and averaging over the transverse photon polarization. That is,

$$\left( \frac{d^5\sigma}{d\mathbf{k}'d\Omega_{\mathbf{k}'}d\Omega_{\mathbf{p}'}} \right)_{\text{lab}} = \frac{1}{2} \sum_{s', \varepsilon} \left( \frac{d^5\sigma(s', \varepsilon)}{d\mathbf{k}'d\Omega_{\mathbf{k}'}d\Omega_{\mathbf{p}'}} \right)_{\text{lab}} . \quad (6)$$

Yet, our prime interest in this work is the calculation of polarization observables: the recoil  $\Lambda$ -polarization ( $\mathcal{P}$ ) and the photon asymmetry ( $\Sigma$ ). The former is defined as [11,12]

$$\mathcal{P} = \sum_{\varepsilon} \left( \frac{d^5\sigma(\uparrow) - d^5\sigma(\downarrow)}{d^5\sigma(\uparrow) + d^5\sigma(\downarrow)} \right)_{\text{lab}} , \quad (7)$$

while the latter by [7,8]

$$\Sigma = \sum_{s'} \left( \frac{d^5\sigma(\perp) - d^5\sigma(\parallel)}{d^5\sigma(\perp) + d^5\sigma(\parallel)} \right)_{\text{lab}} . \quad (8)$$

In these expressions  $\uparrow$  and  $\downarrow$  represent the projection of the spin of the  $\Lambda$ -hyperon with respect to the normal to the scattering plane ( $\mathbf{k}' \times \mathbf{k}$ ), while  $\perp(\parallel)$  represents the out-of-plane(in-plane) polarization of the photon.

### A. Elementary ( $\gamma p \rightarrow K^+ \Lambda$ ) Amplitude

For the elementary photoproduction amplitude we have used a standard model-independent parameterization. The parameterization is given in terms of four Lorentz- and gauge-invariant amplitudes [11,13–15]

$$T(\gamma p \rightarrow K^+ \Lambda) = \sum_{i=1}^4 A_i(s, t) M_i , \quad (9)$$

where the invariant matrices have been defined as:

$$M_1 = -\gamma^5 \not{\varepsilon} \not{k} , \quad (10a)$$

$$M_2 = 2\gamma^5 [(\varepsilon \cdot p)(k \cdot p') - (\varepsilon \cdot p')(k \cdot p)] , \quad (10b)$$

$$M_3 = \gamma^5 [\not{\varepsilon}((k \cdot p) - \not{k}(\varepsilon \cdot p))] \quad (10c)$$

$$M_4 = \gamma^5 [\not{\varepsilon}((k \cdot p') - \not{k}(\varepsilon \cdot p'))] . \quad (10d)$$

Although standard, the above parameterization is not unique. Many other parameterizations — all of them equivalent on shell — are possible [16]. It is our hope that the study of quasifree polarization observables, together with state-of-the-art measurements, will help us elucidate the form of the elementary amplitude.

For the present study, as in all of our earlier photoproduction studies of pseudoscalar mesons [16–18], we transform the standard photoproduction amplitude to a more suitable form. The resultant form is given by

$$T(\gamma p \rightarrow K^+ \Lambda) = F_T^{\alpha\beta} \sigma_{\alpha\beta} + i F_P \gamma_5 + F_A^\alpha \gamma_\alpha \gamma_5 , \quad (11)$$

where tensor, pseudoscalar, and axial-vector amplitudes have been introduced:

$$F_T^{\alpha\beta} = \frac{1}{2} \varepsilon^{\mu\nu\alpha\beta} \varepsilon_\mu k_\nu A_1(s, t) , \quad (12a)$$

$$F_P = -2i [(\varepsilon \cdot p)(k \cdot p') - (\varepsilon \cdot p')(k \cdot p)] A_2(s, t) , \quad (12b)$$

$$F_A^\alpha = [(\varepsilon \cdot p)k^\alpha - (k \cdot p)\varepsilon^\alpha] A_3(s, t) + [(\varepsilon \cdot p')k^\alpha - (k \cdot p')\varepsilon^\alpha] A_4(s, t) . \quad (12c)$$

This form manifests explicitly the Lorentz- and parity-transformation properties of the various bilinear covariants. Note that in the presently-adopted parameterization, no scalar nor vector invariant amplitudes appear. For our calculations we use the hadronic model developed by Williams, Ji, and Cotanch [11]. These authors impose crossing symmetry in their model to develop phenomenologically consistent strong-coupling parameterizations which simultaneously describe the kaon-photoproduction and radiative-capture reactions.

## B. Nuclear-Structure Model

The nuclear-structure model employed here is based on a relativistic mean-field approximation to the Walecka model [5]. For spherically symmetric potentials, such as the strong scalar and vector potentials used here, the eigenstates of the Dirac equation can be classified according to a generalized angular momentum  $\kappa$  [19]. These eigenstates can be expressed in a two component representation; namely,

$$\mathcal{U}_{E\kappa m}(\mathbf{x}) = \frac{1}{x} \begin{bmatrix} g_{E\kappa}(x) \mathcal{Y}_{+\kappa m}(\hat{\mathbf{x}}) \\ i f_{E\kappa}(x) \mathcal{Y}_{-\kappa m}(\hat{\mathbf{x}}) \end{bmatrix} , \quad (13)$$

where the spin-angular functions are defined as:

$$\mathcal{Y}_{\kappa m}(\hat{\mathbf{x}}) \equiv \langle \hat{\mathbf{x}} | l \frac{1}{2} j m \rangle ; \quad j = |\kappa| - \frac{1}{2} ; \quad l = \begin{cases} \kappa , & \text{if } \kappa > 0; \\ -1 - \kappa , & \text{if } \kappa < 0. \end{cases} \quad (14)$$

The Fourier transform of the relativistic bound-state wavefunction — needed to compute the photoproduction cross section [see Eq. (5)] — can now be easily evaluated . We obtain,

$$\mathcal{U}_{E\kappa m}(\mathbf{p}) \equiv \int d\mathbf{x} e^{-i\mathbf{p}\cdot\mathbf{x}} \mathcal{U}_{E\kappa m}(\mathbf{x}) = \frac{4\pi}{p} (-i)^l \begin{bmatrix} g_{E\kappa}(p) \\ f_{E\kappa}(p)(\boldsymbol{\sigma} \cdot \hat{\mathbf{p}}) \end{bmatrix} \mathcal{Y}_{+\kappa m}(\hat{\mathbf{p}}) , \quad (15)$$

where we have written the Fourier transforms of the radial wave functions as

$$g_{E\kappa}(p) = \int_0^\infty dx g_{E\kappa}(x) \hat{j}_l(px) , \quad (16a)$$

$$f_{E\kappa}(p) = (\text{sgn}\kappa) \int_0^\infty dx f_{E\kappa}(x) \hat{j}_{l'}(px) . \quad (16b)$$

Note that in the above expression we have introduced the Riccati-Bessel function in terms of the spherical Bessel function [20]:  $\hat{j}_l(z) = z j_l(z)$  and that  $l'$  is the orbital angular momentum corresponding to  $-\kappa$  [see Eq. (14)].

### C. Closed-form Expression for the Photoproduction Amplitude

Having introduced all relevant quantities, we are now in a position to evaluate the (square of the) photoproduction amplitude [Eq. (5)]. As mentioned earlier, polarization observables are computed within the framework of a relativistic PWIA. That this is, indeed, an excellent approximation seems to be justified by the nonrelativistic analysis presented in Refs. [6–8]. Without distortions, the evaluation of the  $\Lambda$  propagator is now standard due to an algebraic “trick” that appears to be used for the first time by Casimir [21,22]:

$$S(p') \equiv \sum_{s'} \mathcal{U}(\mathbf{p}', s') \bar{\mathcal{U}}(\mathbf{p}', s') = \frac{\not{p}' + M_\Lambda}{2M_\Lambda} ; \quad \left( p'^0 \equiv E_\Lambda(\mathbf{p}') = \sqrt{\mathbf{p}'^2 + M_\Lambda^2} \right) . \quad (17)$$

Subsequently others — Feynman being apparently the first one — used this trick to reduce the “complex” computation of covariant matrix elements to a simple and elegant evaluation of traces of Dirac  $\gamma$ -matrices. These trace-techniques have been used here to compute free polarization observables (note that free polarization observables will serve as the baseline for comparison against bound-nucleon calculations). In principle, one does not expect that these useful trace-techniques will generalize once the nucleon goes off its mass shell. Yet, simple algebraic manipulations — first performed to our knowledge by Gardner and Piekarewicz [9] — show that a trick similar to that of Casimir holds even for bound spinors. Indeed, the validity of their result rests on the following simple identity:

$$\sum_m \mathcal{Y}_{+\kappa m}(\hat{\mathbf{p}}) \mathcal{Y}_{\pm\kappa m}^*(\hat{\mathbf{p}}) = \pm \frac{2j+1}{8\pi} \left\{ \frac{1}{\sigma \cdot \hat{\mathbf{p}}} \right\} , \quad (18)$$

which enables one to introduce the notion of a “bound-state propagator”. That is,

$$\begin{aligned} S_\alpha(\mathbf{p}) &\equiv \frac{1}{2j+1} \sum_m \mathcal{U}_{\alpha,m}(\mathbf{p}) \bar{\mathcal{U}}_{\alpha,m}(\mathbf{p}) \\ &= \left( \frac{2\pi}{p^2} \right) \begin{pmatrix} g_\alpha^2(p) & -g_\alpha(p) f_\alpha(p) \sigma \cdot \hat{\mathbf{p}} \\ +g_\alpha(p) f_\alpha(p) \sigma \cdot \hat{\mathbf{p}} & -f_\alpha^2(p) \end{pmatrix} \\ &= (\not{p}_\alpha + M_\alpha) , \quad (\alpha = \{E, \kappa\}) . \end{aligned} \quad (19)$$

Note that we have defined the above mass-, energy-, and momentum-like quantities as

$$M_\alpha = \left( \frac{\pi}{p^2} \right) [g_\alpha^2(p) - f_\alpha^2(p)] , \quad (20a)$$

$$E_\alpha = \left( \frac{\pi}{p^2} \right) [g_\alpha^2(p) + f_\alpha^2(p)] , \quad (20b)$$

$$\mathbf{p}_\alpha = \left( \frac{\pi}{p^2} \right) [2g_\alpha(p) f_\alpha(p) \hat{\mathbf{p}}] , \quad (20c)$$

which satisfy the “on-shell relation”

$$p_\alpha^2 = E_\alpha^2 - \mathbf{p}_\alpha^2 = M_\alpha^2 . \quad (21)$$

The evident similarity in structure between the free and bound propagators [Eqs. (17) and (19)] results in an enormous simplification; we can now employ the powerful trace techniques

developed by Feynman to evaluate all polarization observables — irrespective if the nucleon is free or bound to a nucleus. It is important to note, however, that this enormous simplification would have been lost if distortion effects would have been incorporated in the propagation of the emitted  $\Lambda$ -hyperon or  $K^+$ -meson.

To provide a feeling for the enormous simplification entailed by the above trick, we derive now an example assuming, for mere simplicity, that the photoproduction amplitude contains only the tensor term [ $M_1$  in Eq. (10)]. For this simplified case the square of the unpolarized photoproduction matrix element [Eq. (5)] becomes proportional to:

$$\begin{aligned} |\mathcal{M}|^2 &\rightarrow |A_1|^2 \left( -\frac{1}{2} g_{\mu\nu} \right) \text{Tr} \left[ \gamma^5 \gamma^\mu \not{k} (\not{p}_\alpha + M_\alpha) \gamma^5 \gamma^\nu \not{k} (\not{p}' + M_\Lambda) \right] \\ &= \frac{1}{2} |A_1|^2 \left[ \text{Tr} \left( \gamma^\mu \not{k} \not{p}_\alpha \gamma_\mu \not{k} \not{p}' \right) - M_\alpha M_\Lambda \text{Tr} \left( \gamma^\mu \not{k} \gamma_\mu \not{k} \right) \right] \\ &= 8 |A_1|^2 (k \cdot p_\alpha) (k \cdot p') . \end{aligned} \tag{22}$$

This result is, indeed, simple and illuminating. Although including the full complexity of the elementary amplitude requires the evaluation of many such terms (not all of them independent) the evaluation of any one of those terms is not much more complicated than the one presented above. Yet, to automate this straightforward but lengthy procedure, we rely on the *FeynCalc 1.0* [23] package with *Mathematica 2.0* to calculate all traces involving  $\gamma$ -matrices. The output from these symbolic manipulations was then fed into a FORTRAN code to obtain the final numerical values for all different polarization observables.

### III. RESULTS AND DISCUSSION

We start the discussion of our results by examining the role of the relativistic dynamics on the polarization observables. On Fig. 2 we display the recoil polarization ( $\mathcal{P}$ ) of the  $\Lambda$ -hyperon and the photon asymmetry ( $\Sigma$ ) as a function of the kaon scattering angle for the knockout of a proton from the  $p^{3/2}$  orbital in  $^{12}\text{C}$ . The polarization observables were evaluated at a photon energy of  $E_\gamma = 1400$  MeV and at a missing momentum of  $p_m = 120$  MeV (this value is close to the maximum in the momentum distribution of the  $p^{3/2}$  orbital; see Fig. 6). Note that in this figure — and all throughout this paper — we compute all observables in the laboratory system using the quasifree condition:  $\omega = \sqrt{q^2 + M_\Lambda^2} - M_N$ . The insensitivity of our results to the relativistic dynamics is evident. Indeed, for the case of the recoil polarization the two curves cannot even be resolved in the figure. We have also examined these effects on the unpolarized cross section and found them insignificant as well. Note that our “nonrelativistic” results were obtained by adopting the free-space relation in the determination of the lower-components of the bound-state wavefunction. This represents our best attempt at reproducing nonrelativistic calculations, which employ free, on-shell spinors to effect the nonrelativistic reduction of the elementary amplitude.

Next we examine the nuclear dependence of the polarization observables. Fig. 3 displays the recoil polarization and the photon asymmetry for the knockout of a valence proton for a variety of nuclei, ranging from  $^4\text{He}$  all the way to  $^{208}\text{Pb}$ . That is, we have computed the knockout from the  $1s^{1/2}$  orbital of  $^4\text{He}$ , the  $1p^{3/2}$  orbital of  $^{12}\text{C}$ , the  $1p^{1/2}$  orbital of  $^{16}\text{O}$ , the  $1d^{3/2}$  orbital of  $^{40}\text{Ca}$ , and the  $3s^{1/2}$  orbital of  $^{208}\text{Pb}$ . We have included also polarization

observables from a single proton to establish a baseline for comparison against our bound-nucleon calculations. The insensitivity of the recoil polarization ( $\mathcal{P}$ ) to the nuclear target is striking, indeed. As soon as the quasifree process takes place from a proton bound to a “lump” of nuclear matter, the recoil polarization becomes insensitive to the fine details of the lump. Moreover, the deviations from the free value (shown with the filled circles) are significant. This indicates important modifications to the elementary amplitude in the nuclear medium and suggests that the recoil polarization of the  $\Lambda$ -hyperon might represent a fundamental property of nuclear matter. For the photon asymmetry ( $\Sigma$ ) the effect, although still significant, does not seem to be as impressive as for the case of the recoil polarization. For example,  $\Sigma$  is no longer independent of the nuclear target. While  $^4\text{He}$  and  $^{12}\text{C}$  are almost identical to the free case,  $^{16}\text{O}$  and  $^{40}\text{Ca}$  — while practically indistinguishable from each other — depart significantly from the free value. Moreover, the photon asymmetry does not seem to saturate, as evince by the  $^{208}\text{Pb}$  results. On the other hand, the shape of the photon asymmetry is preserved in going from the free nucleon value all the way to  $^{208}\text{Pb}$ .

Having established the independence of polarization observables to the nuclear target and to relativistic effects — particularly for the case of the recoil polarization — we are now in a good position to discuss the sensitivity of these observables to the elementary amplitude (note that the insensitivity of polarization observables to final-state interaction has been shown convincingly in Ref. [8]). In the crossing-symmetric model of Ref. [11] two equally good fits to kaon-photoproduction and radiative-capture data were provided. These are the two sets, labeled “C1” and “C2”, that we adopt here. We note that other choices — perhaps more sophisticated and up to date — can be easily incorporated into our calculation.

We display in Fig. 4 the differential cross section as a function of the kaon scattering angle for the knockout of a proton from the  $p^{3/2}$  orbital in  $^{12}\text{C}$ . Again, the photon incident energy and the missing momentum have been fixed at 1400 MeV and 120 MeV, respectively. Although there are noticeable differences between the two sets at small angles, these differences essentially disappear past  $\theta_K \simeq 20^\circ$ . Much more significant, however, are the differences between the two sets for the case of the polarization observables displayed in Fig. 5. The added sensitivity to the choice of amplitude exhibited by the polarization observables should not come as a surprise; unraveling subtle details about the nuclear dynamics is the hallmark of polarization observables. We have compared our results in Figs. 4-5 to the non-relativistic results of Ref. [8] and have found nice agreement between them, particularly in the case of the set C1. Both calculations predict similar behavior for all the observables and differ only in the fine details. Knowing that relativity plays no role in the calculation — and knowing that the polarization observables are very sensitive to the choice of amplitude — the differences between the two sets of calculations must lie solely in the choice of elementary amplitude.

Finally, we display in Fig. 6 the cross section as a function of missing momentum for the  $p^{3/2}$  orbital in  $^{12}\text{C}$  using a different kinematical setting. Here we have kept the photon incident energy fixed at 1400 MeV but have set the momentum transfer  $q$  at 400 MeV. To a large extent the cross sections represents — up to an overall normalization factor — the momentum distribution of the  $p^{3/2}$  orbital. Indeed, the peak in the cross section is located at  $p_m \approx 110$  MeV, which is also the position of the maximum in the momentum distribution. Moreover, to further appreciate the similarities between the two we have included the energy-like parameter  $E_\alpha$ , up to an arbitrary scale. As seen from Eq. (20b),  $E_\alpha$



is directly proportional to the momentum distribution of the bound-proton wavefunction. The similarities between the cross section and the momentum distribution are indisputable. Note that the cross section dies out for  $p_m > 250$  MeV. This region of high-momentum components is sensitive to short-range correlations, which are beyond the scope of our simple mean-field description. Thus, the tail of the photoproduction cross section can be used to test more sophisticated models of nuclear structure.

#### IV. CONCLUSIONS

We have computed polarization observables — the recoil polarization of the  $\Lambda$ -hyperon and the photon asymmetry — for the quasifree  $K^+$  photoproduction reaction from nuclei. Due to the insensitivity of polarization observables to distortion effects, a relativistic plane-wave impulse approximation was developed. For the elementary amplitude we used the relativistic crossing-symmetric model developed by Williams, Ji, and Cotanch in Ref. [11], while for the nuclear structure we employed a relativistic mean-field approximation to the Walecka model [5]. In this manner the quasifree amplitude was evaluated without recourse to a nonrelativistic reduction, as the full relativistic structure of the amplitude was maintained.

By assuming the validity of the relativistic plane-wave impulse approximation an enormous simplification ensued: by introducing the notion of a bound-state propagator — as was done for the first time by Gardner and Piekarewicz in Ref. [9] — the mathematical structure of all quasifree observables was casted in a manner analogous to that of the elementary process. Thus, we brought the full power of Feynman's trace techniques to bear into the problem. We stress that the relativistic formalism presented here can be applied with minor modifications to most quasifree knockout studies, at least in the plane-wave limit.

In addition of being insensitive to distortions effects, we found polarization observables almost independent of the target nucleus and insensitive to relativistic effects. The only effect that polarization observables appear to be sensitive to is the elementary amplitude. As free polarization observables provide a baseline against which possible medium effects may be inferred, we conclude that quasifree polarization observables might be one of the cleanest tools for probing modifications to the elementary amplitude in the nuclear medium. Deviations from their free values are likely to stem from a modification of the elementary interaction inside the nuclear medium due, for example, to a change in resonance parameters. Indeed, for the kinematics adopted in this work ( $E_\gamma = 1.4$  GeV or  $\sqrt{s} \approx 1.9$  GeV) one should be very sensitive to the formation, propagation, and decay of the  $P_{13}(1900)$  and  $F_{17}(1990)$   $N^*$ -resonances [24]. The meson photoproduction (and electroproduction) programs at various experimental facilities — such as TJNAF, NIKHEF, and MAMI — should shed light into the physics of this interesting and fundamental problem.

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# FIGURES

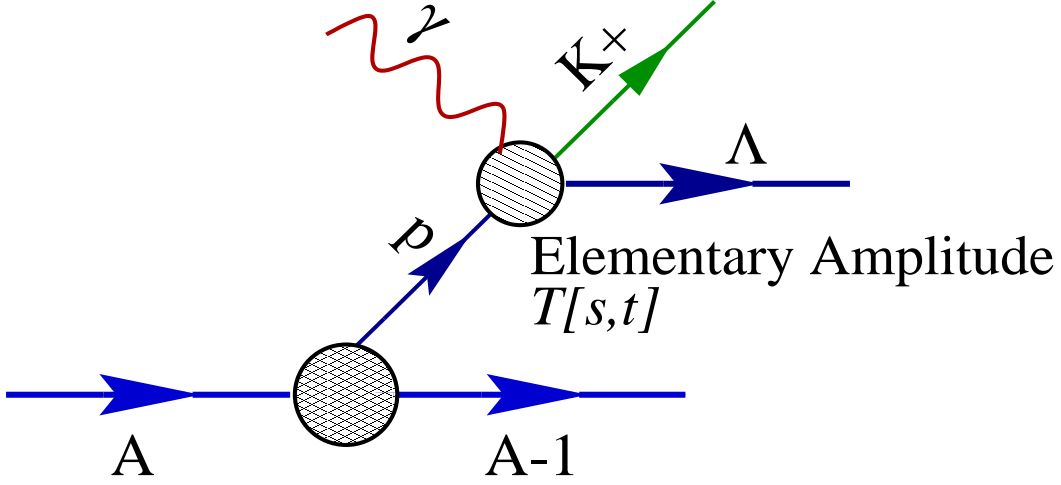


FIG. 1. Representation of the quasifree photoproduction of a  $\Lambda$ -hyperon in a plane-wave impulse-approximation approach.

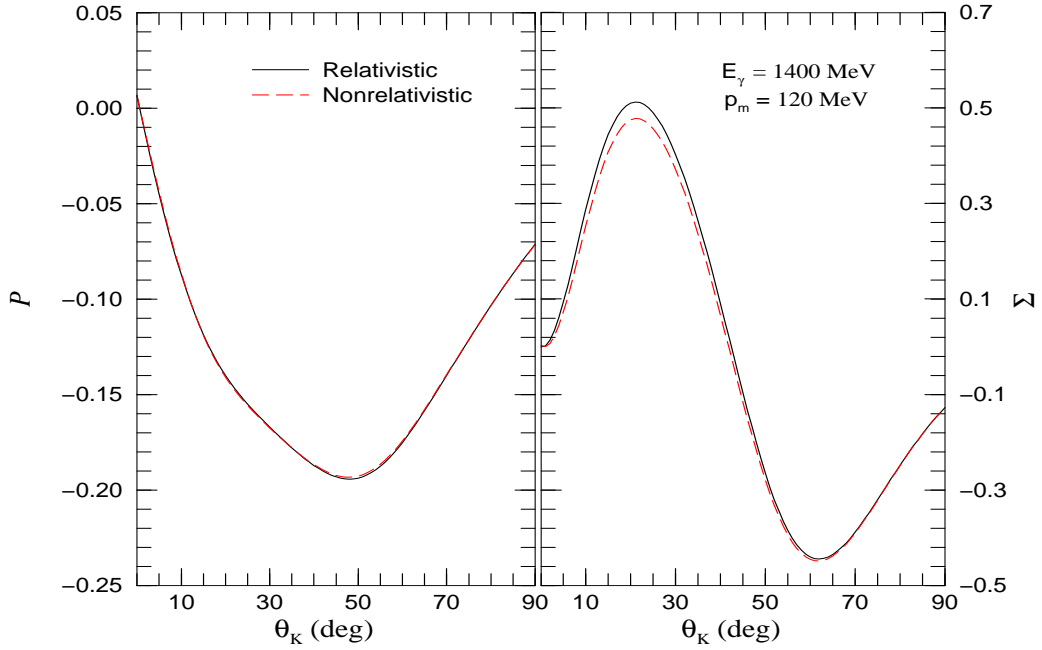


FIG. 2. Comparison between relativistic and nonrelativistic calculations of the recoil polarization of the  $\Lambda$ -hyperon ( $P$ ) and the photon asymmetry ( $\Sigma$ ) as a function of the kaon scattering angle for the knockout of a proton from the  $p^{3/2}$  orbital in  $^{12}\text{C}$ .

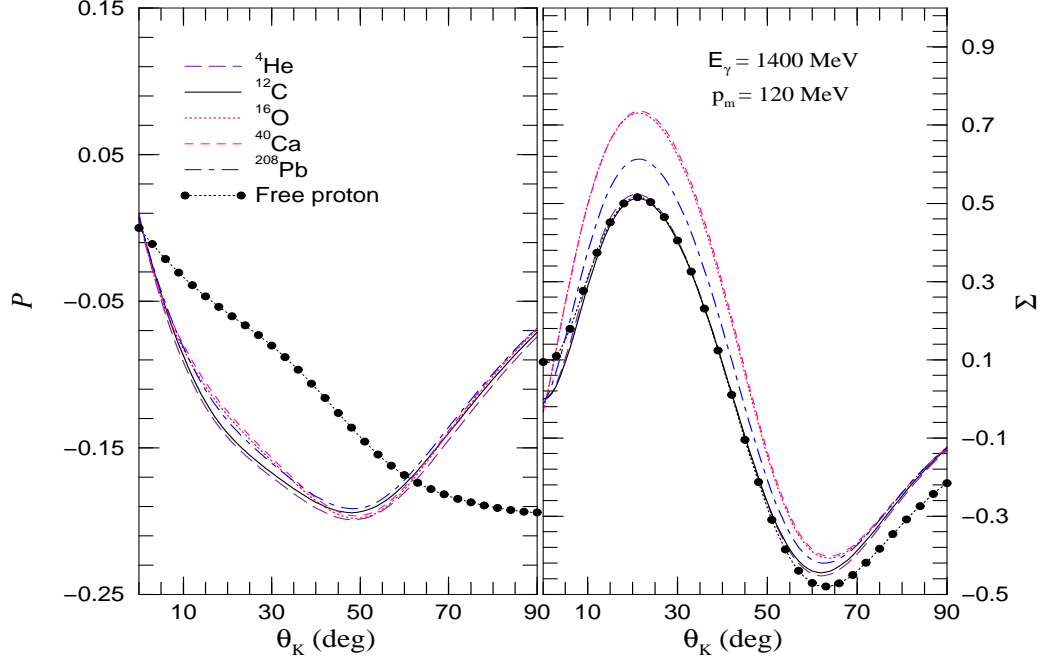


FIG. 3. The recoil polarization of the  $\Lambda$ -hyperon ( $\mathcal{P}$ ) and the photon asymmetry ( $\Sigma$ ) as a function of the kaon scattering angle for the knockout of a valence proton from a variety of nuclei. The photoproduction from a free proton is depicted with the filled circles.

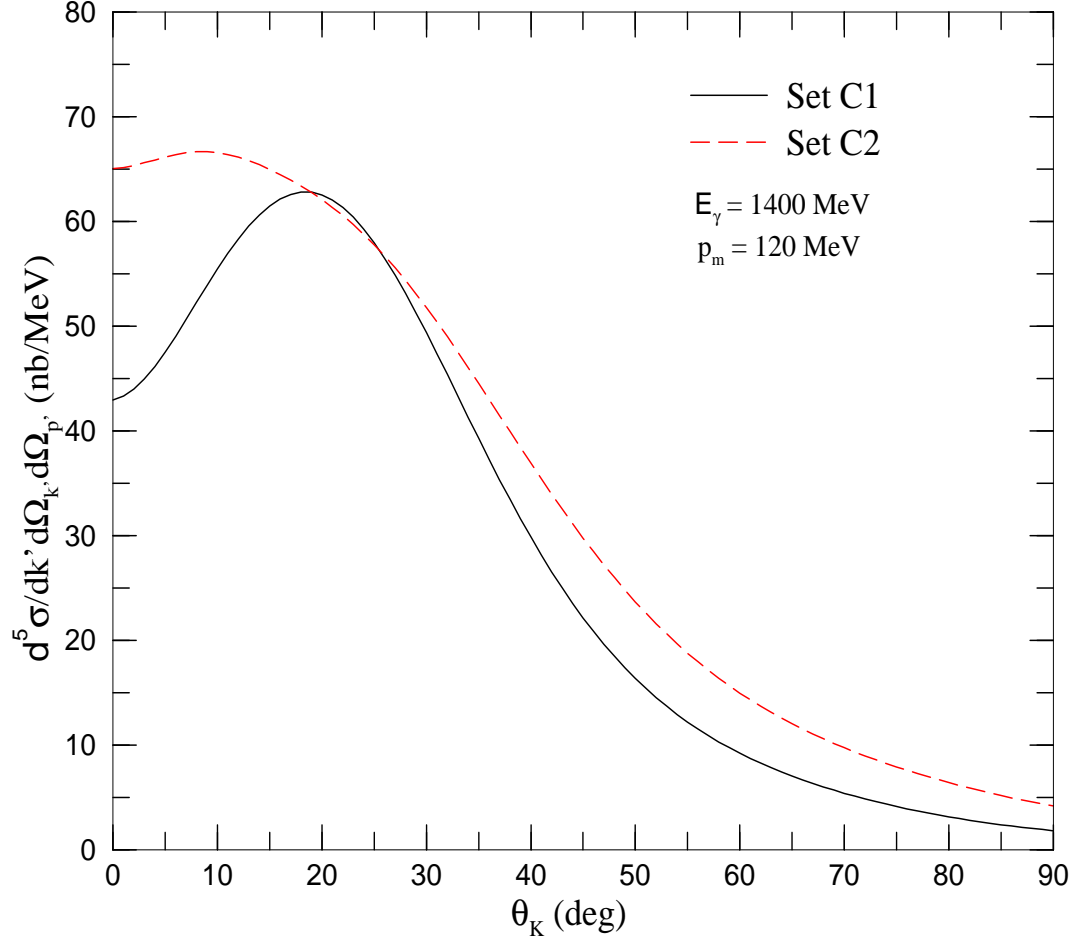


FIG. 4. The differential cross section as a function of the kaon scattering angle for the knockout of a proton from the  $p^{3/2}$  orbital in  $^{12}\text{C}$  using two different sets for the elementary amplitude: C1 and C2.

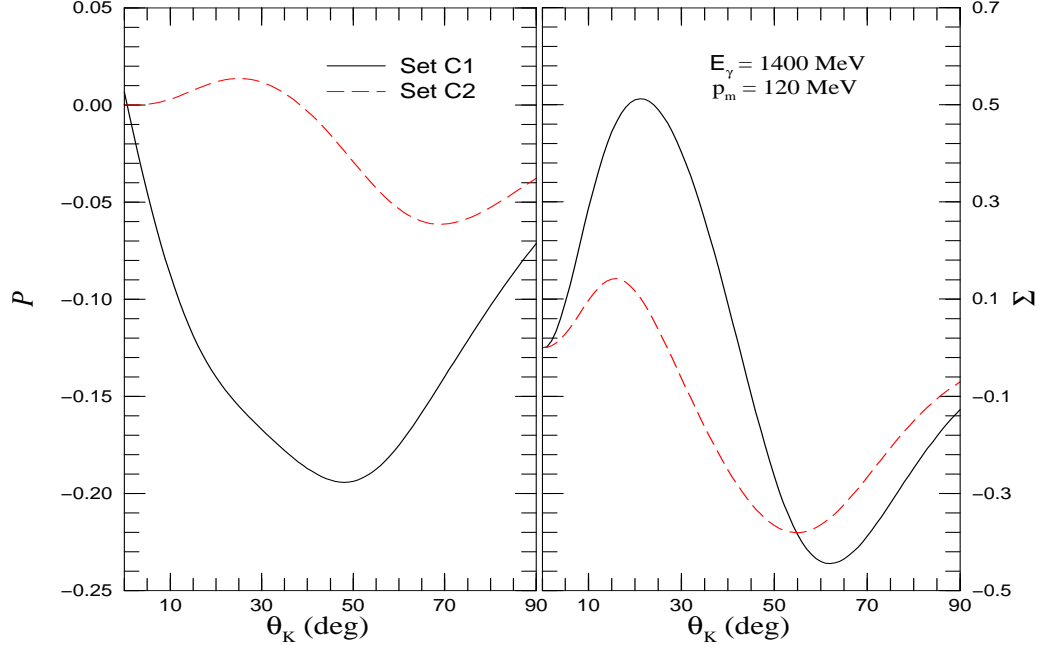


FIG. 5. The recoil polarization of the  $\Lambda$ -hyperon ( $\mathcal{P}$ ) and the photon asymmetry ( $\Sigma$ ) as a function of the kaon scattering angle for the knockout of a proton from the  $p^{3/2}$  orbital in  $^{12}\text{C}$  using two parameterizations of the elementary amplitude: C1 and C2.

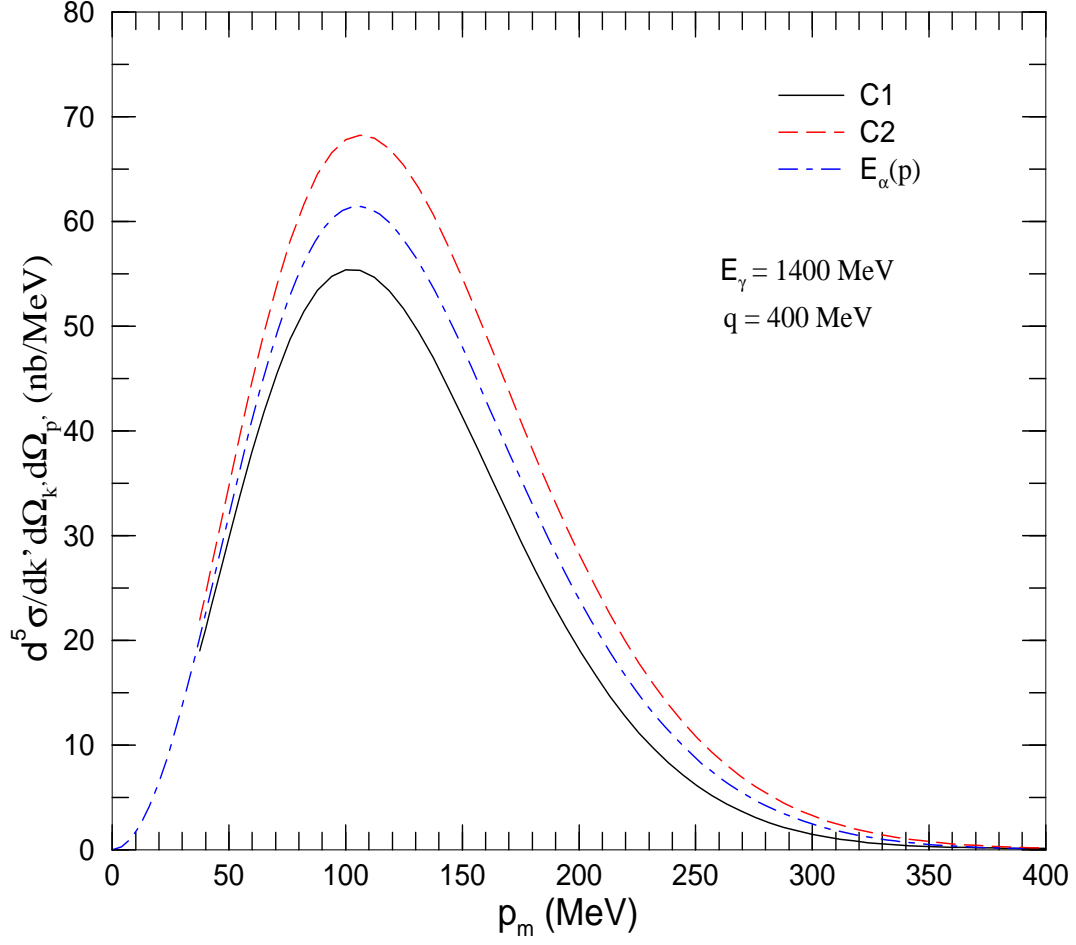


FIG. 6. The differential cross section as a function of the missing momentum for the knockout of a proton from the  $p^{3/2}$  orbital of  $^{12}\text{C}$  using two parameterizations of the elementary amplitude: C1 and C2. The figure includes also the  $E_\alpha$  parameter (up to an arbitrary scale) which is proportional to the momentum distribution of the bound-proton wavefunction (see text for details).